D-Particles, D-Instantons, and A Space-Time Uncertainty Principle in String Theory

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The purpose of this talk is to review some considerations by the present author on the possible role of a simple space-time uncertainty relation toward nonperturbative string theory. We first motivate the space-time uncertainty relation as a simple space-time characterization of the fundamental string theory. In the perturbative string theory, the relation can be regarded as a direct consequence of the world-sheet conformal invariance. We argue that the relation captures some of the important aspects of the short-distance dynamics of D-particles described by the effective super Yang-Mills matrix quantum mechanics, and also that the recently proposed type IIB matrix model can be regarded as the simplest realization of the space-time uncertainty principle.

1 Introduction

1.1 Motivation: Basic principles behind string theory?

Many people working in string theory now believe that the string theory as the prime candidate toward the unified quantum theory of all interactions including gravity has now reached at a very crucial stage. Namely, the recent developments related to the string dualities completely changed the naive view on the theory: The theory is now not even a theory of strings, but rather it suggests some more fundamental degrees of freedom behind strings. The strings turn out to be only one form of the various possible degrees of freedom of the theory which appears in the special perturbative regime of the theory. In fact, varieties of extended objects (called "branes") with various dimensions appear in different regimes of the theory, and they may play equally important roles as the one-dimensional objects, strings. Although a suspicion that the string may not be the fundamental objects of the theory has long been held by many people, this is the first time that we have concrete evidence and handles for that. We may now attack the problem of formulating the string theory nonperturbatively in a well-defined manner.

In spite of these impressive developments, however, we do not yet have any clear understanding on the basic principles which underlie the string theory. In view of this, it seems worthwhile to present some considerations on a new uncertainty principle with respect to space and time [1, 2], which, albeit speculative and hence rather vague yet, explains in a very simple way some of the crucial qualitative aspects of both the perturbative and nonperturbative excitations of the theory and might play some role in the quest of basic principles of string theory.

1.2 Origin of the space-time uncertainty relation

To motivate our discussion below, let us first briefly recall the meaning of the fundamental string scale. Obviously, one of the most significant feature of the string theory is that it introduces the fundamental unit of length, $\sqrt{\alpha'} \equiv \ell_s$. The string length ℓ_s replaces the role of the Planck length in ordinary approaches to quantizing gravity within the framework of local field theory. In string theory, the gravitational interaction is built-in in a completely unified way with all other particle interactions. In the unit in which the 10 dimensional Newton constant and the light velocity are unity, the Planck constant h is related to the string length as

$$h = g_s^2 \ell_s^8 \tag{1}$$

where g_s is the dimensionless string coupling given by the vacuum expectation value of the dilaton $g_s = e^{\phi}$. Note that, in this "gravity unit", the dimension of mass is L^{D-3} in terms of the length dimension L, while the dimension of the Planck constant $\hbar$ is L^{D-2}, in D dimensional space-time. Thus the classical limit $h \to 0$ keeping gravity finite corresponds either to $\ell_s \to 0$ or to $g_s \to 0$. The latter situation means that the universe to be described by the classical physics is composed of classically extended string objects. This does not fit our experience, and we are forced to adopt the former option. In this case, the introduction of the Planck's quantum is directly equivalent to introducing the string scale. From this viewpoint, the string theory should perhaps be interpreted as a theory of the quantized space-time, in which the ordinary space-time continuum plays a somewhat analogous role as the classical phase space in ordinary quantum mechanics.

Let us now try to reinterpret the string theory from this viewpoint. Since in any theory of quantum gravity we expect the appearance of a limiting length beyond which the ordinary concept of the continuum space-time should be invalidated, it is important to identify the manner how such a limitation arises in the fundamental string theory. One such characterization of the theory comes from a reinterpretation of the ordinary time-energy uncertainty relation.

$$\Delta T \Delta E \geq \hbar. \tag{2}$$

Now the energy of a string is roughly proportional to the length of the string X_l, the string tension $T \sim \frac{\hbar}{\alpha'}$ being the proportional constant:

$$\Delta E \sim \frac{\hbar}{\alpha'} \Delta X_l. \tag{3}$$

Thus we have the relation for the uncertainties with respect to the time interval and the spatial length [1, 2],

$$\Delta T \Delta X_l \geq \ell_s^2. \tag{4}$$

It is tempting to assume this relation as a fundamental uncertainty relation which governs the short-distance behavior of the string theory as a "theory of quantized space-time", since it expresses the limitation of the ordinary concept of the contin-

uum space-time directly in terms of the string scale ℓ_s and the space-time area. [a] The dual role of the time and the (longitudinal [b]) spatial lengths is a natural space-time expression of the original Regge and resonance duality which has of course been the original principle, underling the string model (or its old name, "dual resonance model"). The Regge and the resonance behaviors correspond to the regimes, $\Delta X_l \to \infty$ and $\Delta T \to \infty$, respectively.

Let us first consider a high-energy scattering by which we can probe the small time scale $\Delta T \to 0$. If we suppose that the uncertainty relation is nearly saturated in the high energy Regge-pole regime with a large s-channel energy and a small t-channel momentum and that the interaction occurs locally along the longitudinal direction of string as in the light-cone picture for the string, the amplitude is expected to be roughly proportional to $\Delta X_l \sim \ell_s^2/\Delta T \propto E$ which implies [6] that the intercept $\alpha(0)$ of the leading Regge trajectory is 2, from the relation $E \sim E^{\alpha(t)-1}$, as it should be in a theory containing the graviton.

In contrast to this, in the high energy fixed-angle scatterings with large s-and t-channel momenta [4], we are trying to probe the region where both the time and the spatial scales are small. Clearly, such a region is incompatible with the space-time uncertainty relation. The exponential fall-off of the perturbative string amplitudes in this limit may be interpreted as a manifestation of this property. According to the space-time uncertainty relation, at least one of the two length scales must be larger than the string scale ℓ_s.

Another method of probing the small time region $\Delta T \to 0$ is to use black holes. Any finite time interval in the asymptotic space-time region corresponds to infinitely short time interval on the black hole horizons because of the infinite red shift. This means that in any scattering event of a string with a black hole, the spatial extension of the string at the horizon is effectively infinite. Thus we expect that the description of string-black hole scattering should be very different from that of ordinary local field theory. For example, this might imply that as for the observer at infinity no information loss occurs in the scattering with a black hole, as has been suggested by Susskind in the name of "black hole complementarity" [7]. Unfortunately, however, since only reliable discussion of string- black hole scatterings is through the low-energy effective field theory approximation, we have no satisfactory example to see how that difference is and what its significance is. In connection with this, it might be worthwhile to recall a very peculiar behavior of a matrix-model description for the $1+1$-dimensional $SL(2,R)/U(1)$ black hole, as has been proposed in ref. [8].

[a]Historical remark: This viewpoint was first suggested by the present author [1] just a decade ago, but unfortunately has not been developed further at that time except for a later discussion given in [2]. Independently of this suggestion, an intimately related concept of the 'minimal length' was suggested by others [3 4] and has got popularity.

[b]Here it is important to discriminate the length scales in the longitudinal and transverse directions with respect to the string. The length scale appeared in the uncertainty relation (4) is the one along the longitudinal direction. As is well known, the (squared) transverse length scale grows logarithmically with energy used to probe the strings. This explains the linearly rising Regge trajectory in the Regge-pole behavior. Note that in the special case of the strings in 1+1 dimensions, there is a longitudinal size but no transverse size, and, hence, the high-energy limits of the perturbative amplitudes show power-law behaviors, apart from the phase factors associated with the external legs. For a discussion on the physics of the 1+1 dimensinal string scattering, we refer the reader to ref. [5].

Note that the behavior $\Delta X_l \to \infty$, being the extension in the longitudinal length, is valid even in $1+1$ dimensions. Our intuition from local field theory would lose its validity when we consider the string-black hole scattering, especially when it comes to the localization property [9] of the interactions in the sense of target space-times. For a review on some controversial issues related to this problem, I would like to refer the reader to ref. [10].

What is the implication of the space-time uncertainty relation in the large time region $\Delta T \to \infty$? The relation tells us that we may have a small spatial scale $\Delta X_l \to 0$. This is at first sight somewhat contrary to our intuition, since any stationary string states always have an extension of order of ℓ_s. However, we can probe the structure of strings only through the scattering experiments, and therefore only observables are the S-matrix elements. The possibility of the non-linear σ model approach shows that the asymptotic states of any string states can be regarded as the local space-time fields. It seems reasonable to interpret that the appearance of the local fields corresponds to the property $\Delta X_l \to 0$. This, however, does not show the possibility of directly probing the short spatial scales in string theory. The space-time uncertainty relation shows that doing this necessarily requires us to spent long time intervals, $\Delta T \to \infty$, or low energies. As we shall see later, the D-particles indeed provides us such a possibility.

1.3 Conformal invariance and the space-time uncertainty relation

Now before proceeding to discuss the implication of the space-time uncertainty relation for the nonperturbative aspects of the string theory, we briefly describe its connection [2] to the world-sheet conformal invariance. One well known property of string perturbation theory, characterizing the short distance behavior of the string theory which is in a marked contrast to the local field theory, is the modular invariance. This shows that there is a natural ultraviolet cut-off of order ℓ_s^{-1} in string theory. The origin of the modular invariance can be traced back to a simple symmetry of the string amplitude for a world-sheet parallelogram. Consider the Polyakov amplitude for the mapping: The boundary conditions are of Dirichlet type :

$$X^\mu(\tau,0) = X^\mu(\tau,b) = \delta^{\mu 0}\frac{A\tau}{a},$$

$$X^\mu(0,\sigma) = X^\mu(a,\sigma) = \delta^{\mu 1}\frac{B\sigma}{b}.$$

where A, B are the lengths in space-time corresponding to the lengths τ, σ on the Riemann sheet using the conformal metric $g_{ab} = \delta_{ab}$. Then apart from a power behaved pre-factor, the amplitude takes the form

$$\exp[-\frac{1}{\ell_s^2}(\frac{A^2}{\Gamma} + \frac{B^2}{\Gamma^*})] \tag{5}$$

where

$$\Gamma \equiv \frac{a}{b},\ \Gamma^* \equiv \frac{b}{a}, \quad (\Gamma \cdot \Gamma^* = 1). \tag{6}$$

Due to the conformal invariance, the amplitude depends on the Riemann sheet parameters only through the ratio Γ or Γ^*, which are called the extremal length

and the conjugate extremal length, respectively. Clearly, the relation $\Gamma\Gamma^* = 1$ leads to the uncertainty relation [2]

$$\Delta T \Delta X \sim \langle A \rangle \langle B \rangle \sim \ell_s^2,$$

and the symmetry $(A, \Gamma) \leftrightarrow (B, \Gamma^*)$ is the origin of the modular invariance of a torus amplitude. It is only at the critical space-time dimensions where the conformal anomaly cancels that the pre-factor enjoys the required symmetry.

This form of the amplitude is reminiscent of the matrix element for a Gaussian wave packet state $< x|g(\Delta x) >\sim \exp -\frac{1}{2}(\frac{x}{\Delta x})^2$ saturating the Heisenberg' s uncertainty relation $\Delta x \Delta p = \hbar$ in the Wigner representation.

$$\mathcal{O}(x,p) \equiv \int dy\, \mathrm{e}^{\frac{ipy}{\hbar}} < x - \frac{1}{2}y|\, \mathcal{O}_{\Delta X}\, |x + \frac{1}{2}y > \propto \exp -\Big[(\frac{x}{\Delta x})^2 + (\frac{p}{\Delta p})^2\Big], \quad (7)$$

where $\mathcal{O}_{\Delta x} = |g(\Delta x) >< g(\Delta x)|$ is the density operator representing the Gaussian state. The correspondence, $(\lambda, A, B, \sqrt{a/b}, \sqrt{b/a}) \leftrightarrow (\hbar, x, y, \Delta x/\sqrt{\hbar}, \Delta p/\sqrt{\hbar})$, suggests that we integrate over the uncertainty Δx in analogy with the integration over the modular parameter in string theory. Thus the statistical density operator $\rho \equiv \int d(\Delta x)\, \mathcal{O}_{\Delta x}$ is an analogue of the string amplitude.

2 D-particle dynamics and the space-time uncertainty relation

Let us now discuss the possible relevance of our space-time uncertainty relation to the nonperturbative aspect of the string theory. One of the most important observation in the recent development is that the Dirichlet branes are solitonic excitations with Ramond-Ramond charge [11] which are crucial ingredients for the validity of the S-duality of string theory. The dynamics of the D-branes are described by the collective degrees of freedom in terms of open strings with Dirichlet boundary conditions. The collective coordinates are matrices living on the D-branes and associated with the open-string vertex operators. The matrix nature of the collective degrees of freedom comes from the Chan-Paton factor of the open-string vertex operators.

For example, for the D0-branes (or D-"particles") in the type IIA theory, the collective variable is one-dimensional Hermitian matrices $X^a_{ij}(t)$ and their fermionic partners, where t is the time along the world line of the D-particle and the space index a runs from 1 to 9. . The Chan-Paton indices i, j run from 1 to N with N being the number of the D-branes involved. The effective action [12] for the collective fields $X^\mu_{ij}(t)$ is nothing but the dimensional reduction of 10 dimensional $U(N)$ super Yang-Mills theory to 0+1 dimensions at least for sufficiently low energy phenomena:

$$S = \int dt \frac{1}{2g_s} \mathrm{Tr}\Big\{ \dot{X}^a \dot{X}^a - \frac{1}{2}[X^a, X^b]^2 + fermions \Big\}. \quad (8)$$

Here we used the unit $\ell_s = 1$ for simplicity and assume the $A_0 = 0$ gauge. This effective action clearly shows that the system along the classical flat direction $[X^a, X^b] = 0$ corresponds to N free nonrelativistic particles of mass $1/g_s$ whose coordinates are the eigenvalue of the matrices X^a.

Is it possible to probe arbitrary short distances by a scattering experiment using these D-particles? [13,14] The interaction of the D-particles are governed by the off-diagonal elements of the matrices which corresponds to the lowest modes of the open strings stretching between the D-particles. The lowest frequency of the off-diagonal elements $X_{ij}^{(\mu)}$ is of order $|\lambda_i - \lambda_j|$ where the λ_i is the coordinate of the i-th D-particle. This means that the characteristic time scale $\Delta T \sim \frac{1}{|\lambda_i - \lambda_j|}$ increases as we try to probe shorter distances. This is just the property required by the space-time uncertainty relation. We can then expect that a scattering of very slow D-particles would make possible to probe very short distances. However, quantum mechanics tells us that if we spent too long time, the wave packets of the D-particles spread and wash out the information on the short distances. The spreading of the wave packet is smaller if the mass of the particles are heavier. In our case, the mass of the D-particles is proportional to the inverse string coupling $1/g_s$. Thus there must be a characteristic spatial scale which gives the order of the shortest possible distances probed by the D-particle scattering. This is easily obtained by the scaling argument. By redefining the time and the coordinate matrices as

$$X = g_s^{1/3} \tilde{X}, \;\; t = g_s^{-1/3} \tilde{t}, \tag{9}$$

we can completely eliminate the string coupling from the effective action. Thus we see that the characteristic spatial and time scales, restoring the string scale, are given by $g_s^{1/3} \ell_s$ and $g_s^{-1/3} \ell_s$, respectively. As is well known, the spatial scale $g_s^{1/3} \ell_s$ is nothing but the 11 dimensional Planck scale of the M-theory [15]. On the other hand, the time scale is given by the inverse power $g_s^{-1/3} \ell_s$ as is required by the space-time uncertainty relation. We note that for the validity of the above dimensional arguments, the existence of the supersymmetry is important, since if the supersymmetry did not exist the zero-point oscillation would generate a long-range potential for the eigenvalue coordinates and would ruin the simple scaling property of the classical effective action.

These results are readily derived, without explicitly using the effective action, if one first assumes the space-time uncertainty relation as an underlying principle of the string theory including nonperturbative objects [16]. Consider the scattering of two D-particles of mass ℓ_s/g_s with the impact parameter of order ΔX and the relative velocity v which is assumed to be much smaller than the light velocity. Then the characteristic interaction time ΔT is of order $\frac{\Delta X}{v}$. Since the impact parameter is of the same order as the longitudinal length of the open strings mediating the interaction of the D-particles, we can use the space-time uncertainty relation in the form

$$\Delta T \Delta X \sim \ell_s^2 \Rightarrow \frac{(\Delta X)^2}{v} \sim \ell_s^2$$

This gives the order of the magnitude for the possible distances probed by the D-particle scatterings with velocity $v \, (\ll 1)$.

$$\Delta X \sim \sqrt{v} \ell_s. \tag{10}$$

The same result is obtained from the effective low-energy action by using the Born-Oppenheimer approximation [14] for the coupling between the diagonal and

off-diagonal matrix elements. The spreading of the wave packet of the D-particle during the time interval $\Delta T \sim \frac{\ell_s}{\sqrt{v}}$ is easily estimated as

$$\Delta X_w \sim \Delta T \Delta_w v \sim \frac{g_s}{v} \ell_s$$

where $\Delta_w v \sim g_s \frac{1}{\sqrt{v}}$ is the uncertainty of velocity caused by the ordinary uncertainty relation for a nonrelativistic particle with energy uncertainty, $\Delta E \sim 1/\Delta T \sim \frac{\sqrt{v}}{\ell_s} \sim \frac{v \Delta_w v \ell_s}{g_s}$. The condition $\Delta X_w (= g_s \frac{\ell_s}{v}) < \Delta X (= \sqrt{v} \ell_s)$ leads

$$\Delta X > g_s^{1/3} \ell_s. \tag{11}$$

Actually, by considering the systems with many D-branes we can probe the smaller spatial scales than the 11 dimensional Planck length. For example, when we consider a D-particle in the presence of many ($=N$) coincident D4-branes, the effective kinetic term for the D-particle is given by [14]

$$S_{\text{eff}} = \int dt [\frac{1}{2g_s \ell_s} (1 + \frac{N g_s \ell_s^3}{r^3}) v^2 + \mathcal{O}(N v^4 \ell_s^6 / r^7)] \tag{12}$$

where r is the distance between the D-particle and the D4-branes. When the distance r is much shorter than $(N g_s)^{1/3} \ell_s$, the effective mass of the D-particle is given by $m \sim N \ell_s^2 / r^3 \gg \frac{1}{g_s \ell_s}$. We can then easily check that the spread of the D-particle wave packet can be neglected during the time $t \sim v^{-1/2} \ell_s$, compared with the scale determined above from the consideration of the space-time uncertainty relation $\sqrt{v} \ell_s$ for large N. This allows us to probe arbitrary short lengths with respect to the distance between the D-particle and D4-branes and hence to see even the singular nature of the effective metric in the action (12). Since the time scale grows indefinitely, however, we cannot talk about the interaction time in any meaningful way in the limit of short spatial distance.

These considerations indicate, I believe, the universal nature of the space-time uncertainty relation as one of the possible principles underlying the nonperturbative string theory.

As is well known now, an exciting interpretation for the effective action (8) is to regard it as the *exact* action of the M-theory for the 9 (=11-2) transverse degrees of freedom in the infinite-momentum frame in 11 dimensions (so called 'M(atrix) theory') [17]. This ansatz has already passed a number of nontrivial tests. In the 11 dimensional M-theory which is believed to be the underlying theory of the type IIA string theory, the D0-brane is interpreted as the Kaluza-Klein excitation of the 11 dimensional massless fields. The compactification radius of the 11th dimension is $g_s \ell_s$, and hence the unit of the Kaluza-Klein momentum in 11 dimensions is just the mass of D0-branes. The fundamental assumption is that the momentum of the D0-brane in the 11th direction is always equal to this single unit of the momentum and hence the total number N of the D0-branes is just proportional to the total 11th momentum. Thus going to the infinite momentum frame is equivalent to the large N limit. Combined our discussion with this proposal, it seems natural to try to construct possible covariant formulations of the M-theory by taking account

the space-time uncertainty relation. My own attempt toward this direction has so far been unsuccessful. The covariant formulation necessarily requires all of the space-time coordinates as matrices. The crucial step would thus be to identify a natural higher symmetry such that it allows us to go to the light-cone gauge effective action (8) in which the light-cone time is reduced to a single variable and the total 11th momentum, measured by the basic unit $1/g_s$, becomes equal to N. Hence, necessarily, the order of the matrices itself becomes a dynamical variable. It turns out that this is very nontrivial to implement in the framework of the matrix models. From physical side, one of the difficulty is that in the ordinary Lorentz frame, we have perhaps to treat D-particles and anti-D-particles simultaneously. However, at present, almost nothing is known concerning about the brane-anti-brane interactions, except that the usual perturbative description in terms of the open strings breaks down near the string scale [18]. In order to achieve the covariant formulation, we would have to make some acrobatic twists to the ordinary matrix model.

Remarks

1. The above scaling argument can be extended to Dp-branes. The longitudinal (including the time direction) and transverse lengths always scale oppositely: The powers are $g_s^{-1/(3-p)}$ and $g_s^{1/(3-p)}$, respectively. This reflects the fact that the interaction of D-branes are mediated by open strings satisfying the space-time uncertainty relation for the time and the longitudinal lengths of the open string. Note that the longitudinal direction of the open strings corresponds to the transverse directions for the D-branes. For example, for $p = 1$ (D-string), the characteristic spatial scale is $g_s^{1/2}\ell_s$. The meaning of this scale is not yet fully understood. In any case (provided $p \geq 0$), we see that the space-time uncertainty relation is satisfied, although the characteristic spatial and time scales vary for each specific case. The case $p = -1$ (D-instanton) is very special in the sense that all the space-time directions are 'transverse' and there is no time evolution; but we shall argue, in the next section, that this case is also consistent with the space-time uncertainty relation. [c]

2. On the other hand, if we go to the strong coupling regime of the string theory, we expect that the D-branes become the lighter excitation modes than the fundamental strings. In particular, in the type IIB theory, the conjectured S-duality relation says that the role of the fundamental strings and the D-strings are interchanged in a completely symmetrical way. Then the main interaction among the D-strings should be through the ordinary splitting and merging of the D-strings. Since the string scale is invariant under the S-duality transformation $g_s \leftrightarrow 1/g_s$, we expect that our space-time uncertainty relation is satisfied even in the strong coupling limit of the type IIB theory. Thus it is plausible that the uncertainly relation is universally valid independently of the

[c] As for a possible interpretation of the string-D-instanton interactions [19] from the viewpoint of the space-time uncertainty relation, see [16]. Essentially, the instanton-string interaction is necessarily in the regime $\Delta T \to 0$, and correspodingly is governed by the massless long-range exchange interactions $\Delta X \to \infty$.

strength of the string coupling, although it is tested only in the weak coupling string theory. It should however be expected that in the intermediate coupling regions, it would be difficult to saturate the uncertainty relations since in this region the halo around the D-branes (F-strings) consisting of the cloud of the F-strings (D-brane) should be generally important around the string scale.

3 Type-IIB matrix models as a possible realization of the space-time uncertainty relation

3.1 D-instanton and the type IIB matrix model

The D0-brane, D-particle, is the lowest dimensional excitation of the type IIA superstring theory which allows only even dimensional D-branes. The type IIB theory, on the other hand, allows only odd dimensional D-branes. In particular, the object of the lowest dimension is the D-1-brane, namely, the D-instanton with $p = -1$. For the D-instanton, all of the space-time directions, including even the time direction, becomes the transverse directions. The interaction of the D-instantons are therefore described by the open strings which obeys the Dirichlet boundary condition with respect to all the directions. Hence the effective action for the D-instantons, at least for low-energy phenomena, are described by the 10 dimensional super Yang-Mills theory reduced to a space-time point:

$$S_{\rm eff} = -\frac{1}{g_s}\mathrm{Tr}([X_\mu, X_\nu])^2 + \text{fermionic part} \tag{13}$$

which is covariant from the beginning. We again suppress the string scale by taking the unit $\ell_s = 1$. The D-instanton is rather special in the sense that the only "physical" degrees of freedom are the lowest massless modes of the open strings. The reason is that only on-shell states satisfying the Virasoro condition are the zero-momentum "discrete" states of the Yang-Mills gauge fields which are identified with X_μ's and their super partners. Other infinite number of the massive degrees of freedom of the open strings are excluded as the solution of the Virasoro condition and therefore should be regarded as a sort of auxiliary fields. We then expect that the exact effective action for them should be expressible solely in terms of the collective fields which are nothing but the lowest modes X_μ. If we further assume that all the higher dimensional objects including the fundamental strings could be various kinds of bound states or collective modes composed of infinitely many D-instantons, [d]we may proceed to postulate that a matrix model defined at a single point might give a possible nonperturbative definition of the type IIB string theory. We however note that this is a very bold assumption: For example, unlike the case of the M(atrix) theory, here there is no rationale why we can neglect the configuration in which both instantons and anti-instantons are present. If we can only hope that suitable choice of the model might take into account such effect implicitly.

A concrete proposal along this line was indeed put forward first by Ishibashi, Kawai, Kitazawa and Tsuchiya [20]. Their proposal is essentially to regard the above

[d]For example, in usual gauge field theories, a monopole can formally be interpreted as being composed of infinite number of instantons along a line.

simple action as the exact action in the following form.

$$Z_{\rm IKKT} = \sum_N \Big(\prod_\mu \int d^{N^2}X_\mu\Big) d^{16}\psi \exp S[X,\psi,\alpha,\beta] \tag{14}$$

$$S[X,\psi,\alpha,\beta] \equiv -\alpha N + \beta \mathrm{Tr}_N \frac{1}{2}[X_\mu, Y_\nu]^2 - \cdots \tag{15}$$

with $\alpha \propto \frac{1}{g_s \ell_s^4}, \beta \propto \frac{1}{g_s}$ and the summation over N is assumed. Indeed, it has been shown that the model reproduces the long-distance behavior of D-brane interactions expected from the low-energy effective theory, the type IIB supergravity in 10 dimensions. Very recently, they have presented further arguments [21] which suggest that the model can reproduce the light-cone string field theory.

3.2 On the Schild action formulation of string theory

Let us now discuss the connection of the space-time uncertainty relation with the type IIB models of this type. Ishibashi et al motivated their proposal by the formulation of Green-Schwarz superstring action in the Schild gauge [22]. In fact, the formulation of the string theory in the Schild gauge [22] is very natural from our viewpoint. So, let me begin from a discussion of the Schild action.

$$S_2 = -\int d^2\xi\, e\Big\{\frac{1}{e^2}[-\frac{1}{2\lambda^2}(\epsilon^{ab}\partial_a X^\mu \partial_b X^\nu)^2] + 1\Big\} + \cdots \tag{16}$$

where $e = e(\xi)$ is an auxiliary field with the same transformation property with the volume density of the world sheet and $\lambda = 4\pi\alpha'$. The classical equation of motion for the string coordinate takes the same form as that from the Nambu-Goto action which is obtained from S_2 by eliminating the auxiliary field e. Alternatively, we can also derive the Polyakov action directly as follows.

The action can be made quadratic by introducing three auxiliary fields $t^{12} = t^{21}, t^{11}, t^{22}$, which forms a tensor density of weight two, as

$$\begin{aligned} S_2 &= -\int d^2\xi\, \frac{1}{\lambda^2 e}[\dot{X}^2 \acute{X}^2 - (\dot{X}\cdot\acute{X})^2] - \int d^2\xi\, e \\ &\equiv S_(t,e,X) - \int d^2\xi \frac{1}{e} \det \tilde{t} \end{aligned} \tag{17}$$

where

$$S(t,e,X) = \int d^2\xi\, \frac{1}{e}[\det t + \frac{1}{\lambda} t^{ab}\partial_a X \cdot \partial_b X] - \int d^2\xi\, e, \tag{18}$$

and we have defined the shifted auxiliary fields $\tilde{t}^{ab}$ by

$$\tilde{t}^{11} = t^{11} + \frac{1}{\lambda}\acute{X}^2, \quad \tilde{t}^{22} \equiv \tilde{t}^{21} = t^{22} + \frac{1}{\lambda}\dot{X}^2, \quad \tilde{t}^{12} = t^{12} - \frac{1}{\lambda}\dot{X}\cdot\acute{X}. \tag{19}$$

The new action $S(t,e,X)$ containing the auxiliary field t^{ab} is quantum-mechanically equivalent to the original Schild action since the difference is a free quadratic form

$\int d^2\xi \frac{1}{e} \det \tilde{t}$ with respect to the auxiliary fields $\tilde{t}^{ab}$, apart from an induced ultra-local measure which we included in the definition of the measure for e :

$$Z_S \equiv \int \frac{[dX][de]}{[d(diff_2)]}\, \mathrm{e}^{-S_2} \propto \int \frac{[dX][de][dt]}{[d(diff_2)]}\, \mathrm{e}^{-S(t,e,X)}. \tag{20}$$

Let us then make a change of variables

$$t^{ab} \to g^{ab}, e \to \tilde{e}$$

where g_{ab} transforms as the standard world-sheet metric and $\tilde{e}$ as a scalar, $t^{ab} = g^{ab}e^2, e = \tilde{e}\sqrt{g}$, and get

$$Z_S \propto \int \frac{[d\tilde{e}][dg][dX]}{[d(diff_2)]} \exp\Big(-\int d^2\xi(\tilde{e}^3 - \tilde{e})\sqrt{g} - \frac{1}{\lambda}\int d^2\xi \tilde{e}\sqrt{g}g^{ab}\partial_a X \cdot \partial_b X\Big). \tag{21}$$

At this point, we can assume that the measure in the partition function Z_S is defined such that the total integration measure $[d\tilde{e}][dX][dg]$ obtained in these transformations is reparametrization invariant. We then decompose the measure $[dg]$ as usual

$$\frac{[dg]}{[d(diff_2)]} = [d\mathrm{Weyl}]\frac{[dg]}{[d\mathrm{Weyl}][d(diff_2)]}.$$

The functional metrics for the Weyl modes and auxiliary scalar fields $\tilde{e}$ are, in the conformal gauge $g_{ab} = \delta_{ab}\,\mathrm{e}^{\phi}$ for simplicity,

$$\|\delta(\mathrm{e}^{\phi})\|^2 = \int d^2\xi \mathrm{e}^{-\phi}(\delta \mathrm{e}^{\phi})^2, \tag{22}$$

$$\|\delta\tilde{e}\|^2 = \int d^2\xi \mathrm{e}^{\phi}(\delta\tilde{e})^2, \tag{23}$$

respectively. Note that in the former expression we have chosen the integration variable to be the density itself $\mathrm{e}^{\phi} = \sqrt{g}$. Now, when the conformal anomaly cancels, there is no induced kinetic term arising from the integration measure $[dg][dX]/[d\mathrm{Weyl}][d(diff_2)]$ for the Weyl mode in the world-sheet metric g_{ab} which therefore completely decouples from the second term of the action in the expression (21). Then, the integration over the Weyl mode gives a δ-function constraint $\tilde{e} = 1$.

Here we assumed that the integration measure in the partition function is chosen such that the diffeomorphism invariant functional δ-function formula is valid,

$$\int [d\tilde{e}][d\sqrt{g}]f[\tilde{e}]\exp(\int d^2\xi\sqrt{g}\tilde{e}) = f[0],$$

for arbitrary $f[\tilde{e}]$ which is independent of $\sqrt{g}$ as a functional of a scalar field $\tilde{e}$. Formally, this amounts to postulating that the integration measure is defined such that the factors $\mathrm{e}^{-\phi}$, e^{ϕ} in the functional metrics (22) and (23), respectively, cancel each other in going from the original functional metrics (22) and (23) to $\overline{\|\delta(\mathrm{e}^{\phi})\|}^2 \equiv \int d^2\xi(\delta(\mathrm{e}^{\phi}))^2$ and $\overline{\|\delta\tilde{e}\|}^2 \equiv \int d^2\xi(\delta\tilde{e})^2$, and that the integration range for $\sqrt{g} \equiv \mathrm{e}^{\phi}$ is in fact extended to the whole imaginary axis. The first assumption

can be justified by adopting the point-splitting regularization scheme for the formal Jacobian $J = |\det(\mathrm{e}^{\phi(\xi_1)-\phi(\xi_1')+\phi(\xi_2)-\phi(\xi_2')}\delta(\xi_1,\xi_1';\xi_2,\xi_2')|$ associated with the transition $[d\tilde{e}][d\sqrt{g}] = \overline{[d\tilde{e}][d\sqrt{g}]}\,J$, such that the sets of the discretized world-sheet points arising from the two factors $\mathrm{e}^{-\phi}$, e^{ϕ} are the same, $\{\xi_i\} = \{\xi_i'\}$. The δ-function $\delta(\xi_1,\xi_1';\xi_2,\xi_2')$ can be regularized by, e.g., the reparametrization invariant heat kernel method. The assumption for the integration range is not inconsistent since the fields g_{ab} come from the auxiliary fields t^{ab}. Also we have discarded the nonsensical solution $\tilde{e} = 0$ which can be excluded by making a justifiable assumption that the measure for $\tilde{e}$ vanishes at $\tilde{e} = 0$. Note that the Gaussian integration over the shifted field $\tilde{t}^{ab}$ in (17) indeed produces such a factor.

Thus the final result is, apart from a proportional numerical constant, that

$$Z_S = \int \frac{[dX][dg]}{[d(diff_2)][d\mathrm{Weyl}]}\,\mathrm{e}^{-S_P^e} \tag{24}$$

in critical space-time dimensions where S_P^e is the Euclidean Polyakov action. Obviously, the argument goes through for superstrings by including necessary fermionic variables once the additional terms for the bosonic Schild action to ensure the supersymmetry are determined. Although the above argument is not rigorous, it seems now fairly clear that we have the quantum string theory using the Schild action, which is equivalent to the standard formulation based on the Polyakov action.

3.3 The space-time uncertainty relation form the Schild action

The Schild action has no manifest Weyl (or conformal) invariance. Nonetheless, it is equivalent with the Polyakov action when the conformal anomaly cancels. Where is then the conformal structure hidden? It is easy to see that the Virasoro constraint for the canonical variables

$$\mathcal{P}^2 + \frac{1}{4\pi\alpha'}\acute{X}^2 = 0, \tag{25}$$

$$\mathcal{P}\cdot\acute{X} = 0. \tag{26}$$

is obtained after using the constraint coming from the variation with respect to the auxiliary field:

$$\frac{1}{e}\sqrt{-\frac{1}{2}(\epsilon^{ab}\partial_a X^\mu \partial_b X^\nu)^2} = \lambda. \tag{27}$$

Since the Virasoro condition expresses the conformal invariance of the world-sheet field theory, we can say that the conformal invariance of the string theory is encoded in the condition (27) in the Schild-gauge formulation. For this reason, we call the condition (27) the "conformal constraint". On the other hand, we have already emphasized that our space-time uncertainty relation can be regarded as a direct consequence of the conformal invariance. Thus it is natural to expect that the constraint (27) is closely related to the spacetime uncertainty relation. Indeed, using the known relation [23] between the commutator for general Hermitian $N\times N$ matrices and the Poisson bracket in two-dimensional phase space (τ,σ) in the large N limit, we can make the following correspondence,

$$\{X_\mu, X_\nu\} \equiv \frac{1}{e}\epsilon^{ab}\partial_a X_\mu \partial_b X_\mu \leftrightarrow [X_\mu, X_\nu]. \tag{28}$$

$$\mathrm{Tr}\,(\cdots) \leftrightarrow \int d\tau d\sigma e\ (\cdots). \tag{29}$$

Then the constraint coming from the Schild action is interpreted as the one for the commutator of the matrices

$$-\frac{1}{2}([X^\mu, X^\nu])^2 = \lambda^2 I \tag{30}$$

This leads to the inequality in the Minkowski metric

$$\langle([X^0, X^i])^2\rangle\ \geq \lambda^2, \tag{31}$$

which is just consistent with the uncertainty relation for the time and spatial lengths. This seems to be the simplest possible expression of the space-time uncertainty relation in a way compatible with Lorentz covariance [24].

3.4 Microcanonical matrix model as a realization of the space-time uncertainty relation

It is now reasonable to try to construct some matrix model which takes into account the condition (30). In the following, we discuss a provisional model [24] which leads to the model (15) at least for low-energies.

We require that the condition (30) in a weaker form, namely, as an average

$$\langle\frac{1}{2}([X_\mu, X_\nu])^2\rangle \equiv \lim_{N\to\infty}\frac{1}{N}\mathrm{Tr}\frac{1}{2}([X_\mu, X_\nu])^2 = -\lambda^2. \tag{32}$$

Here, $\langle\cdot\rangle$ denotes the expectation value with respect to the $U(N)$ trace as indicated, and the large N limit is assumed to include the case of arbitrary number of D-instantons. The large N limit is also necessary in order to include, for instance, a static configuration of the D-string with the relation

$$[X_0, X_i] \propto I$$

for some spatial direction i. The fundamental partition function is then defined as

$$Z = \int\Big(\prod_{\mu=1}^{10} d^{\infty^2}X_\mu\Big)\mathcal{J}[X]\delta(\langle\frac{1}{2}([X_\mu, X_\nu])^2\rangle + \lambda^2) \tag{33}$$

where $d^{\infty^2}X_\mu$ is the large N limit of the standard $U(N)$ invariant Haar measure and the factor $\mathcal{J}[X]$ is an additional measure factor to be determined below.

Obviously, our model is similar to a microcanonical partition function of classical statistical mechanics. In the latter case, we adopt the Liouville measure of the phase space, appealing to the Liouville theorem and ergodic theory. In our case, based on the interpretation of the instanton collective coordinates, we require the cluster property that the effective dynamics for clusters of instantons which are separated far apart from each other should be independent. The configuration of separated clusters are represented by diagonal blocks of smaller matrices $Y_\mu^{(a)}$ $(N_a \times N_a)$ and $Y_\mu^{(b)}$ $(N_b \times N_b)$ embedded in the original large $N \times N$ matrices

X_μ and the distance between two clusters are measured by the the difference of the center of mass coordinates $\ell_{a,b} \equiv |\frac{1}{N_a}\mathrm{Tr}_a Y_\mu^{(a)} - \frac{1}{N_b}\mathrm{Tr}_b Y_\mu^{(b)}|$. In general, the fluctuations of the off-diagonal elements generate effective long range interaction of the form $\log \ell_{a,b}$ coming from the Vandelmonde measure of the matrix integral. This necessarily violates the cluster property. A natural way out of this problem is to require supersymmetry for the partition function by introducing the fermionic partner to the collective coordinates to cancel the Vandelmonde measure. Of course, if one tries to derive the effective theory for D-instantons from the type IIB string theory, one would have such fermionic partners, automatically. In this way, we are led to introduce the ansatz for the measure factor $\mathcal{J}[X]$

$$\mathcal{J}[X] = \int d^{16}\psi' \exp(\frac{1}{2}\langle\overline{\psi}\Gamma_\mu[X_\mu,\psi]\rangle) \tag{34}$$

where ψ is the hermitian $N\times N$ matrix whose elements are Majorana-Weyl spinors in 10 dimensions and the prime in the integration volume denotes that possible fermion zero-modes should be removed for the partition function. The supersymmetry is easily established after rewriting the partition function by introducing an auxiliary constant multiplier c,

$$Z = \int dc \Big(\prod_\mu d^{\infty^2}X_\mu\Big) d^{16}\psi' \exp\Big[c(\langle\frac{1}{2}([X_\mu,X_\nu])^2\rangle + \lambda^2) + \frac{1}{2}\langle\overline{\psi}\Gamma_\mu[X_\mu,\psi]\rangle\Big]. \tag{35}$$

The action in this expression is invariant under the two (global) supersymmetry transformations

$$\begin{aligned}
\delta_\epsilon\psi &= ic[X_\mu,X_\nu]\Gamma_{\mu\nu}\epsilon && (36)\\
\delta_\epsilon X_\mu &= i\overline{\epsilon}\Gamma_\mu\psi && (37)\\
\delta_\epsilon c &= 0 && (38)\\
\delta_\eta\psi &= \eta && (39)\\
\delta_\eta X_\mu &= 0 && (40)\\
\delta_\eta c &= 0 && (41)
\end{aligned}$$

where the Grassmann spinorial parameters ϵ,η are "global", i.e., the unit matrix with respect to the $U(N)$ indices. These symmetries can also be derived by the similarity with the Green-Schwarz formulation [25] of superstrings with a special gauge condition for the local κ symmetry. The presence of the supersymmetry ensures that the classical configurations which preserves a part of the supersymmetry can enjoy the BPS property. For example, if the commutator $[X_\mu,X_\nu]$ is proportional to the unit matrix, there remains just a half of the supersymmetry.

We emphasize that our argument is still provisional. For instance, we cannot exclude the possibility of lower dimensional space-times, since we can construct the similar models as the dimension reductions of any supersymmetric Yang-Mills theories. However, the 10 dimensions is the largest dimension where the above construction works and only in this dimension the correct long range forces for D-strings emerges.

Now from our viewpoint, the IKKT model can be naturally derived as the effective low-energy theory of many distant clusters of D-branes. We introduce the block-diagonal matrices whose entries are the smaller $N_a \times N_a$ matrices $Y_\mu^{(a)}$ and require the conditions

$$-\sum_{a=1}^{n} \mathrm{Tr}_a \frac{1}{2}[Y_\mu^{(a)}, Y_\nu^{(a)}]^2 = N\lambda^2, \tag{42}$$

$$\sum_{a=1}^{n} N_a = N \ (\to \infty). \tag{43}$$

Now let us suppose to evaluate the fluctuations around the backgrounds by setting $X_\mu = Y_\mu^b + \tilde{X}_\mu$. In the one-loop approximation, the calculation is entirely the same[e] as the IKKT model [20], and it is easy to see that the leading order behavior of the interaction contained in the one-loop effective action for the backgrounds decreases as $O(\frac{1}{\ell_{a,b}^8})$ in the limit of large separation. Remember that here again the supersymmetry is crucial. The distant D-brane systems therefore can be treated as independent objects in this approximation, and hence we can take into account the conditions (42), (43) in a statistical way by introducing two Lagrange multipliers α and β. Then the effective partition function Z_{eff} for the D-brane subsystems described Y_μ within the semi-classical approximation is given by the grand canonical partition function,

$$Z_{\mathrm{eff}} = \sum_N \Big(\prod_\mu \int d^{N^2} Y_\mu\Big) d^{16}\psi \exp S[Y, \psi, \alpha, \beta], \tag{44}$$

$$S[Y, \psi, \alpha, \beta] \equiv -\alpha N + \beta \mathrm{Tr}_N \frac{1}{2}[Y_\mu, Y_\nu]^2 - \frac{1}{2}\mathrm{Tr}\, \overline{\psi}[\Gamma_\mu, Y_\mu]\psi, \tag{45}$$

where by N we denote the order of the background submatrix Y and Tr_N is the corresponding trace. This form is identical with the IKKT model and explains the origin of the two parameters. In our case, however, we have the condition

$$\begin{aligned} -\langle\langle \frac{1}{2}[Y_\mu, Y_\nu]^2 \rangle\rangle &\equiv -\frac{\sum_N \Big(\prod_\mu \int d^{N^2} Y_\mu\Big) d^{16}\psi \, \frac{1}{N}\mathrm{Tr}_N \frac{1}{2}[Y_\mu, Y_\nu]^2 \exp S[Y, \psi, \alpha, \beta]}{\sum_N \Big(\prod_\mu \int d^{N^2} Y_\mu\Big) d^{16}\psi \exp S[Y, \psi, \alpha, \beta]} \\ &= \lambda^2 \end{aligned} \tag{46}$$

ensuring the quantum constraint by which the Lagrange multipliers should in principle be determined. The original microcanonical model has two parameters, string scale $\sqrt{\lambda}$ and N which corresponds to two parameters α, β of the low-energy effective action. In the large N limit, we expect there remains a dimensionless parameter other than the fundamental length. It should be related to the coupling constant of the effective string theory. In the nonperturbative string theory, we expect that the coupling constant should be determined dynamically from the vacuum expectation

[e]Note that the one-loop calculation is reduced to the determinant in both models.

values of the scalar background fields such as dilaton and/or its dual partner, scalar axion. It would be very interesting to see whether the condition (46) is consistent with this expectation. For example, the form (46) can be interpreted as the effective dilaton equation of motion

$$\frac{\partial}{\partial \phi} Z = 0$$

if $\alpha \sim \frac{1}{g_s \ell_s^4}, \beta \sim \frac{1}{g_s}$. As is discussed in ref. [20], this is just the correct parameter dependence required to reproduce the classical D-string interactions.

4 Discussions

There are many reasons that the discussions given here for the relevance and formulation of the space-time uncertainty principle for the string theory are yet provisional and should be regarded only as a motivation for further exploration of the basic principles behind and beyond the string theory.

Among others, one of the most crucial puzzle seems to me why the gravity is contained in the matrix model formulation of string theory. Both in the type IIA and IIB models, we start from the flat space-time background and there is, at least apparently, no symmetry which can be connected to general coordinate transformation and its supersymmetric generalization. Of course, there exists the space-time supersymmetry which is maximal in each case. However, it is not at all clear how it must be elevated to local symmetry. It is possible that such local symmetries are hidden in some form in the the $SU(N)$ symmetries in the large N limit. In that case, our usual understanding of the background space-time should be reconsidered for the matrix-model approaches. We need some correspondence principle which relates the classical geometric field theory to matrix models containing gravity.

Another very important question is how to unify the entirely different realizations of the space-time uncertainty relation in the D-particle (type IIA) and D-instanton (type IIB) cases. In the D-particle case, the time variable is treated just as the ordinary time in quantum mechanics. The space-time uncertainty relation is a consequence of the special form of the Hamiltonian governing the D-particle dynamics. The nature of the space-time uncertainty relation is thus similar to that of the ordinary time-energy uncertainty relation, reflecting the first derivation of the former as a reinterpretation of the latter. On the other hand, in the case of the D-instantons, the time direction is treated as a matrix variable which is of equal footing as the spatial directions. The space-time uncertainty relation is taken into account by imposing the δ-function constraint as a quantum condition. I expect that something which unifies both approaches and has some higher symmetry structure than the $SU(N)$ of the effective actions in order to allow us the light-cone gauge fixing is necessary for constructing the covariant (11 dimensional) formulation of the M-theory in terms of a matrix model. It seems that this requires us to invent a totally new way of formulating dynamics, as is already expressed earlier in this talk.

To conclude, I have discussed the relevance and possible formulations of a simple space-time uncertainty principle. Although I think the discussions provide some sound evidence, there remains much to be further clarified. Also it is very important

to try to combine our considerations with other proposals which are related to black-hole physics in string theory, like e.g., 'black hole complementarity' [7], 'holographic principle' [26], 'correspondence principle for black holes and strings' [27]. The black hole is expected to be one of the most fruitful testing grounds for the nonperturbative string theory.

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